Soft Electromagnetic Radiations From Equilibrating Quark-Gluon Plasma

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We evaluate the bremsstrahlung production of low mass dileptons and soft photons from equilibrating and transversely expanding quark gluon plasma which may be created in the wake of relativistic heavy ion collisions. We use initial conditions obtained from the self screened parton cascade model. We consider a boost invariant longitudinal and cylindrically symmetric transverse expansion of the parton plasma and find that for low mass dileptons ($M \leq 0.3$ GeV) and soft photons ($p_T \leq 0.5$ GeV), the bremsstrahlung contribution is rather large compared to annihilation process at both RHIC and LHC energies. We also find an increase by a factor of 15-20 in the low mass dileptons and soft photons yield as one goes from RHIC to LHC energies.

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I. INTRODUCTION

The study of ultra-relativistic heavy ion collision explores an opportunity to verify the possible occurrence of a phase transition from hadronic matter to deconfined quark matter [1]. In order to detect this exotic phenomenon, various signals have been proposed. Among them, photons and dileptons have long been considered as excellent probes of the early stages of relativistic heavy ion collision. They are the only signals emitted by the Quark Gluon phase that reach the detector directly without being affected by final state interaction. As a result, their detection would provide valuable and reliable information about the moment of their formation.

One would like to understand several aspects of the matter produced in heavy ion collisions, viz., how does the initial partonic system evolve and how quickly does it attain kinetic equilibrium? How quickly, if at all, does it attain chemical equilibrium? And finally, how can we uncover the history of this evolution by studying the spectra of the produced particles, many of which may decouple from the interacting system only towards the end? These and related questions have been actively debated in recent times. Thus, it is believed by now that the large initial parton density may force many collisions among the partons in a very short time and lead to a kinetic equilibrium [2]. The question of chemical equilibration [3] is more involved, as it depends on the time available to the system. If the time available is too short (3–5 fm/c), as at the energies ($\sqrt{s} \simeq 200$ GeV/nucleon) accessible at the Relativistic Heavy Ion Collider (RHIC), the QGP will end its journey far away from chemical equilibrium. At the energies ($\sqrt{s} \simeq 5.5$ TeV/nucleon) that will be achieved at the CERN Large Hadron Collider (LHC) this time could be large (more than 10 fm/c), driving the system very close to chemical equilibration [4–8], if there is only a longitudinal expansion. However, this time is also large enough to enable a rarefaction wave from the surface of the plasma to propagate to the center ($\tau_s \sim R_T/c_s$; R_T is the transverse nuclear dimension and c_s is the speed of sound), thus introducing large transverse velocity (gradients) in the system. The large transverse velocities may impede the chemical equilibration by introducing a faster cooling. The large velocity gradients may drive the system away from chemical equilibration by introducing an additional source of depletion of the number of partons [9] in a given fluid element.

Dileptons and photons are however produced at every stage of the collision and in an expanding system their number is obtained by an integration over the four-volume of the interaction zone. At very early times the temperature is rather high and we have a copious production of large transverse momentum photons, and large mass dileptons. This should give us a reliable information about the initial stages of the plasma. However, the transverse flow of the system is moderate at early times, and their production will only be marginally affected by the flow. By the time the flow and the other aspects of QGP develop, the temperature would have dropped considerably and we have a large production of low transverse momentum photons and low mass dileptons as well. However they will have, in addition to the contributions from the plasma, a large contribution from hadronic reactions. We shall be required to understand them in detail before we can confidently embark upon the task of deciphering the more involved development of the QGP, as it cools (see, e.g.,Ref, [9]).

The soft photons and low mass dileptons, mostly pro-

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duced at later stages when the system is expected to be affected strongly by transverse flow, can provide reliable and useful information about the later stage of the interaction zone. As a consequence, this soft region has been suggested as a possible window for obtaining imprints about the flow. Obviously, the precise location of this window is determined by the initial conditions of the plasma and the dynamics of the space time evolution. We may also add that recent data from CERES experiment [10] at CERN SPS for the $S + Au$ system show a strong excess in the dielectron yield compared to $p + Be$ and $p+Au$ systems below the ρ mass. This result has already been motivated a lot of systematic studies in this mass region [11]. It is generally accepted that bremsstrahlung processes mainly contribute at low masses. In Ref. [12] it has been shown that the description of the CERES data will improve to some extent, when one adds hadronic bremsstrahlung contribution.

For a *fully equilibrated plasma*, the production of low mass dileptons and soft photons due to bremsstrahlung process has been studied in great detail in the literature [13–15]. However, in the recently formulated self-screened parton cascade (SSPC) model [3] early hard scattering produce a medium which screens the long range color fields associated with softer interactions. When two heavy nuclei collide at sufficiently high energy, the screening occurs on the length scale where perturbative QCD still applies. This approach yields predictions for the initial conditions of the forming QGP without the need of any *ad-hoc* virtuality and momentum cut-off parameters. These calculation also show that the QGP likely to be formed in RHIC and LHC energies could be hot and initially far away from chemical equilibrium. With passage of time, chemical reactions among the partons will push it towards chemical equilibration initially, but the large transverse velocity gradient drives the system away from chemical equilibration [9]. It is quite evident that signals which are emitted at QGP phase would somehow reflect the state of non-equilibration at the time of the phase transition.

In the present work, we mainly concentrate on a region of low mass dileptons ($M \leq 0.5$ GeV) and soft photon ($p_T \leq 0.5$ GeV) production due to bremsstrahlung in a more complete manner, *i.e.*, in a chemically equilibrating and transversely expanding quark-gluon plasma. Since the energies of the electromagnetic radiations are low enough, we obtain the production of low mass dileptons and soft photons from bremsstrahlung processes within a soft photon approximation [16]. The goal of this approximation is to separate the process of electromagnetic radiation from the strong interaction components. One also neglects the four momentum of photon in the δ-function, which could be compensated by inserting a phase-space correction factor (see below).

It is worth noting here that in a recent calculation Aurenche et al. [17] have also calculated bremsstrahlung of soft photons (real and virtual) in the frame work of effective perturbative expansion based on resummation of hard thermal loops (HTL). They demonstrated that bremsstrahlung appears only on the two-loop level when the exchange gluon is space like, and its contribution is of same order as that of one-loop. Since the exchanged gluon is space-like, it manifests itself in the rate to the square of the thermal gluon mass ($m_g \sim gT$), in contrast to the one-loop result where only the thermal quark mass appears [18]. For virtual photons this result is of the same order in coupling constant, but it can differ quantatively because of the gluon mass. On the other hand, for real photon bremsstrahlung a factor $(1/g^2)$ arises due to collinear singularities associated with fermionic propagator caused by the vanishing photon mass in the external line of the vacuum polarization diagrams in the two loops. This is compensated by the factor m_g^2 ($\sim g^2T^2$) which leaves the two-loop bremsstrahlung contribution for real soft photon at the same order in gT as that of one-loop. We shall come back to this aspect later.

The paper is organized in the following way. In Sec.II, we briefly recall hydrodynamic and chemical evolution of the plasma in a transverse direction. In Sec.III we present the formulation for bremsstrahlung production of soft photon and low mass dileptons in a chemically equilibrating and transversely expanding plasma. We discuss our results in Sec.IV. Finally in Sec.V, we give a brief summary.

II. HYDRODYNAMIC EXPANSION, AND CHEMICAL EQUILIBRATION

We assume that kinetic equilibrium has been achieved when the momenta of partons become locally isotropic. At the collider energies it has been estimated that, $\tau_i \approx 0.2 - 0.3$ fm/c [2]. Beyond this point, further expansion is described by hydrodynamic equations and the chemical equilibration is governed by a set of master equations which are driven by the two-body reactions ($gg \leftrightarrow q\bar{q}$) and gluon multiplication and its inverse process, gluon fusion ($gg \leftrightarrow ggg$). The other (elastic) scatterings help maintain thermal equilibrium. The hot matter continues to expand and cools due to expansion and chemical equilibration. We shall somewhat arbitrarily terminate the evolution once the energy density reaches some critical value (here taken as $\epsilon_f = 1.45$ GeV/fm^3 [19]).

The expansion of the system is described by the equation for conservation of energy and momentum of an ideal fluid:

$$\partial_\mu T^{\mu\nu} = 0 \,, \qquad T^{\mu\nu} = (\epsilon + P)u^\mu u^\nu + P g^{\mu\nu} \,, \qquad (1)$$

where ϵ is the energy density and P is the pressure measured in the frame comoving with the fluid. The four-velocity vector u^μ of the fluid satisfies the constraint $u^2 = -1$. For a partially equilibrating plasma the distribution functions for gluons and quarks are assumed

$$n_i(E_i, \lambda_i) = \lambda_i \tilde{n}_i(E_i) \quad , \qquad (2)$$

where $\tilde{n}_i(E_i) = (e^{\beta E_i} \mp 1)^{-1}$ is the BE (FD) distributions for gluons (quarks), and λ_i is the fugacity for parton species i, which describes the deviation from chemical equilibrium. This fugacity factor takes into account undersaturation of parton phase space density, *i.e.*, $0 \leq \lambda_i \leq 1$. The equation of state for a partially equilibrated plasma of massless particles can be written as [4]

$$\epsilon = 3P = [a_2 \lambda_g + b_2 (\lambda_q + \lambda_{\bar{q}})] T^4 , \qquad (3)$$

where $a_2 = 8\pi^2/15$, $b_2 = 7\pi^2 N_f/40$, N_f is the number of dynamical quark flavors. Now, the density of an equilibrating partonic system can be written as

$$n_g = \lambda_g \tilde{n}_g, \qquad n_q = \lambda_q \tilde{n}_q, \qquad (4)$$

where $\tilde{n}_k$ is the equilibrium density for the parton species k:

$$\tilde{n}_g = \frac{16}{\pi^2}\zeta(3) T^3 = a_1 T^3, \qquad (5)$$

$$\tilde{n}_q = \frac{9}{2\pi^2}\zeta(3) N_f T^3 = b_1 T^3. \qquad (6)$$

We further assume that $\lambda_q = \lambda_{\bar{q}}$. The equation of state (3) implies the speed of sound $c_s = 1/\sqrt{3}$. We solve the hydrodynamic equations (1) with the assumption that the system undergoes a boost invariant longitudinal expansion along the z-axis and a cylindrically symmetric transverse expansion [20]. It is then sufficient to solve the problem for $z = 0$, because of the assumption of boost invariance.

The master equations [4] for the dominant chemical reactions $gg \leftrightarrow ggg$ and $gg \leftrightarrow q\bar{q}$ are

$$\begin{aligned} \partial_\mu (n_g u^\mu) &= n_g (R_{2\to 3} - R_{3\to 2}) - (n_g R_{g\to q} - n_q R_{q \to g}) , \\ \partial_\mu (n_q u^\mu) &= \partial_\mu (n_{\bar{q}} u^\mu) = n_g R_{g\to q} - n_q R_{q\to g}, \end{aligned} \qquad (7)$$

in an obvious notation. In case of transverse expansion, the master equations can be shown [9] to lead to partial differential equations:

$$\begin{aligned} &\frac{\gamma}{\lambda_g}\partial_t \lambda_g + \frac{\gamma v_r}{\lambda_g}\partial_r \lambda_g + \frac{1}{T^3}\partial_t(\gamma T^3) + \frac{v_r}{T^3}\partial_r(\gamma T^3) \\ &\quad + \gamma \partial_r v_r + \gamma \left(\frac{v_r}{r} + \frac{1}{t} \right) \\ &\quad = R_3 (1 - \lambda_g) - 2 R_2 \left(1 - \frac{\lambda_q \lambda_{\bar{q}}}{\lambda_g^2} \right) , \\ &\frac{\gamma}{\lambda_q}\partial_t \lambda_q + \frac{\gamma v_r}{\lambda_q}\partial_r \lambda_q + \frac{1}{T^3}\partial_t(\gamma T^3) + \frac{v_r}{T^3}\partial_r(\gamma T^3) \\ &\quad + \gamma \partial_r v_r + \gamma \left(\frac{v_r}{r} + \frac{1}{t} \right) \\ &\quad = R_2 \frac{a_1}{b_1} \left(\frac{\lambda_g}{\lambda_q} - \frac{\lambda_{\bar{q}}}{\lambda_g} \right) , \end{aligned} \qquad (8)$$

where v_r is the transverse velocity and $\gamma = 1/\sqrt{1 - v_r^2}$. The R_2 and R_3 related to the rates appearing in (7) are given by,

$$\begin{aligned} R_2 &\approx 0.24 N_f \alpha_s^2 \lambda_g T \ln(1.65/\alpha_s \lambda_g), \\ R_3 &= 1.2 \alpha_s^2 T (2\lambda_g - \lambda_g^2)^{1/2}, \end{aligned} \qquad (9)$$

where the color Debye screening and the Landau - Pomeranchuk - Migdal effect suppressing the induced gluon radiation have been taken into account, explicitly.

The hydrodynamic equations (1) are solved numerically, with the initial conditions obtained from the SSPC model [3], to get $\epsilon(r,t)$ and $v_r(r,t)$, which serve as input into the equations (8) for the fugacities (for details see Ref. [9]). For the convenience of reader, we plot the different parton fugacities in Figs.(1-2) for a transversely expanding plasma likely to be produced in RHIC and LHC collider energies. In the figures N denotes the time.

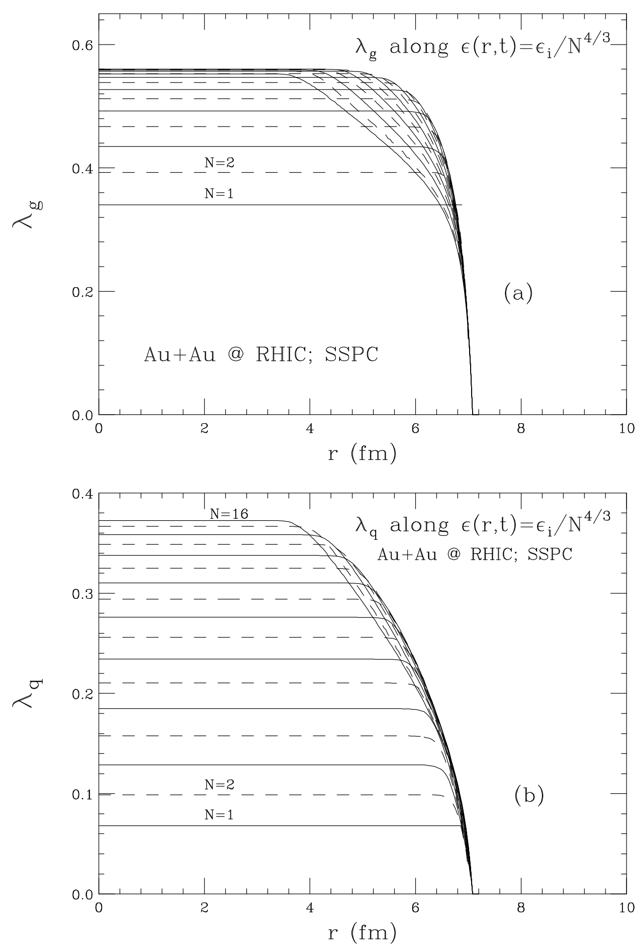

FIG. 1. (a) Gluon fugacities and (b) Quark fugacities along the constant energy density contours for RHIC energies. The constant energy density contours are the solution [9] of the hydrodynamic equation (1) with the initial conditions obtained from SSPC model.

One can see that the quark fugacities always lag the gluon fugacities (see Figs.1 and 2). This is obvious as

SSPC model predicts the gluon dominated plasma. At both the energies the fugacities initially increase with time (denoted by N), then start decreasing because by then the transverse expansion has set in (v_r is non-zero). Thus the region where the transverse velocity is large (later time), the plasma ends it journey farther away from chemical equilibrium. The reason is the velocity gradient causes a depletion in the partons number in the fluid element, which could not be made up by the number of partons produced due to parton chemistry [9], resulting in a decrease of the parton fugacities. The chemical equilibration cannot be achieved at RHIC as the life time (3-4 fm/c) is too small. Due to the longer life time of the plasma phase at LHC energies (12 fm/c), the effects of the transverse expansion are more dramatic, and the entire matter participate in the flow. There is a possibility of approaching chemical equilibration at LHC energies, and goes far away from it as velocity gradient is substantial.

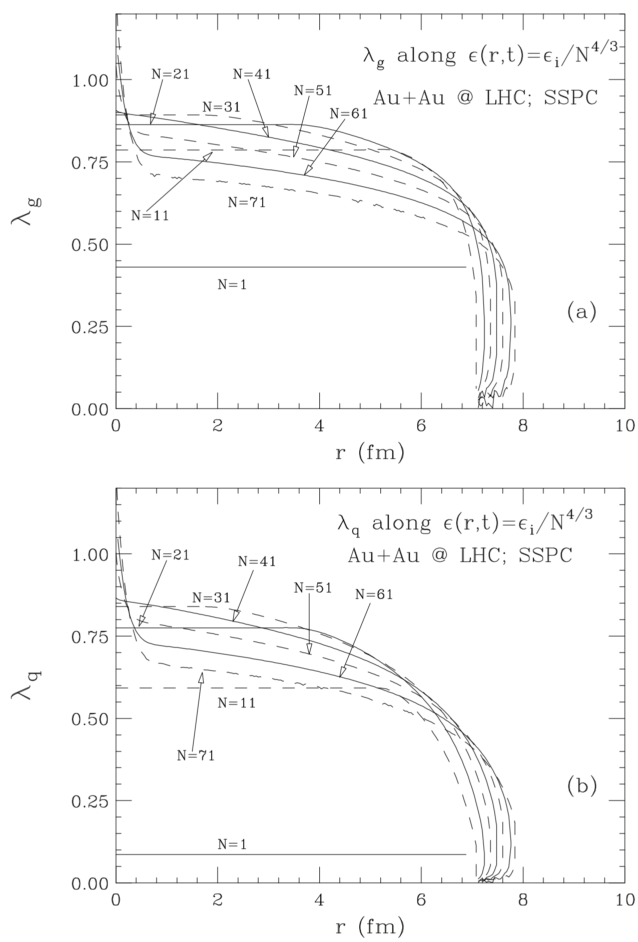

FIG. 2. Same as before, at LHC.

III. PHOTON AND DILEPTON PRODUCTION FROM CHEMICALLY EQUILIBRATING PLASMA

In this section we wish to outline the rate of production of real and virtual photon due to bremsstrahlung in an chemically equilibrating plasma from the following reactions:

$$a \ + \ b \rightarrow c \ + \ d \ + \ \gamma(\gamma^\star) \ . \qquad (10)$$

One can write the matrix element for this processes as

$$\mathcal{M} \ = \ e\mathcal{M}_0 J^\mu \epsilon_\mu, \qquad (11)$$

where $\mathcal{M}_0$ is the matrix element for $ab \rightarrow cd$ and J^μ is the real(virtual) photon current [13,14]. Using kinetic theory one can write down the rate of production of a photon per unit time per unit volume at a temperature T as

$$q_0 \frac{dN}{d^4x \ d^3q} = g_{ab} \frac{1}{2(2\pi)^3} \int \prod_{i=a,b} \left(\frac{d^3p_i}{(2\pi)^3 \ 2E_i} n_i(E_i, \lambda_i) \right) \times \int \prod_{j=c,d} \left(\frac{d^3p_j}{(2\pi)^3 \ 2E_j} \left(1 \pm n_j(E_j, \lambda_j)\right) \right) \times (2\pi)^4 \delta(p_a + p_b - p_c - p_d - q) |\mathcal{M}|^2, \qquad (12)$$

where $g_{ab} = N_a N_b (2s_a + 1)(2s_b + 1)$ is the color and spin appropriate for the reactions in (10). We have also considered the enhancement (suppression) in the exit channel corresponding to boson (fermion). The non-equilibrium distribution function is given in (2). The product of non-equilibrium distribution functions appearing in (12) can be decomposed in terms of equilibrium distributions as follows

$$n_a n_b (1 \pm n_c)(1 \pm n_d) = (1 \pm \lambda_c \tilde{n}_c) \big[\lambda_a \lambda_b \lambda_d \left(1 \pm \tilde{n}_d\right) \tilde{n}_a \tilde{n}_b + \lambda_a \lambda_b \left(1 - \lambda_d\right) \tilde{n}_a \tilde{n}_b \big] \ . \qquad (13)$$

As discussed earlier we neglect the q in the δ-function in the above (12). To compensate this one needs to insert a phase space correction factor Φ, defined as the ratio of three-body phase space to two-body phase space. This can be obtained [12,14,21] as

$$\Phi\left(s, s_2, m_a^2, m_b^2\right) = \frac{s \xi^{1/2}\left(s_2, m_a^2, m_b^2\right)}{s_2 \xi^{1/2}\left(s, m_a^2, m_b^2\right)} \ , \qquad (14)$$

where $\xi(x,y,z) = x^2 - 2x(y+z) + (y-z)^2$, $s_2 = s - 2q_0\sqrt{s}$ for real photon and $s_2 = s + M^2 - 2q_0\sqrt{s}$ for dileptons in the case of equal mass scattering ($m_a = m_c$ and $m_b = m_d$).

Taking into account the non-equilibrium nature of the plasma, the phase space correction, and also inserting $\int ds \ \delta(s - (p_a + p_b)^2)$, the dilepton transverse mass spectrum yield due to bremsstrahlung of a virtual photon ($ab \rightarrow cd\gamma^\star \rightarrow cd\, l^+ l^-$) in the quark phase is obtained from (12) as

$$\frac{dN}{dM^2 d^2M_T dy} = \int \tau\, d\tau\, r\, dr\, d\phi\, d\eta\, \frac{T^6 g_{ab}}{16\pi^4}$$
$$\times \int_{z_{min}}^{\infty} dz \frac{\xi(z^2T^2, m_a^2, m_b^2)}{T^4}\, (1 \pm \lambda_c \tilde{n}_c)$$
$$\times \Big[\lambda_a\lambda_b\lambda_d \sum_{n=0} \left(\pm e^{\beta E_c}\right)^n \frac{\mathcal{K}_1\left(z(n+1)\right)}{n+1}$$
$$+ \lambda_a\lambda_b\left(1-\lambda_d\right)\mathcal{K}_1(z)\Big]\Phi(s, s_2, m_a^2, m_b^2)$$
$$\times \left(E\frac{d\sigma^{l^+l^-}_{ab\to cd}}{dM^2 d^3q}\right), \tag{15}$$

where $z_{min} = (m_a + m_b + M)/T$, $z = \sqrt{s}/T$ and $E_c = \sqrt{p_c^2 + m_a^2}$ with $|\vec{p}_c| = \xi^{1/2}(s, m_a^2, m_b^2)/(2\sqrt{s})$. The cross-section for the process $ab \to cd\, l^+l^-$ is given by

$$E\frac{d\sigma^{l^+l^-}_{ab\to cd}}{dM^2 d^3q} = \frac{d\sigma^{l^+l^-}_{ab\to cd}}{dM^2 d^2M_T dy} = \frac{\alpha^2}{12\pi^3 M^2}\, \frac{\widehat{\sigma}(s)}{M_T^2\, \cosh^2 y}, \tag{16}$$

where M_T is the transverse mass of the dilepton and y is its rapidity, so that, $q_0 = M_T\, \cosh y$. The $\widehat{\sigma}(s)$ is defined as

$$\widehat{\sigma}(s) = \int_{-\xi(s,m_a^2,m_b^2)/s}^{0} dt\, \frac{d\sigma_{ab\to cd}}{dt}$$
$$\times \left(q_0^2 |\epsilon\cdot J|^2_{ab\to cd}\right)\ . \tag{17}$$

The corresponding result for the transverse mass distribution of dileptons from the equilibrating plasma, due to the annihilation process, is given [9] by

$$\frac{dN_{l^+l^-}}{dM^2\, d^2M_T\, dy} = \frac{\alpha^2}{2\pi^3}\lambda_q\lambda_{\bar{q}} e_q^2 \int \tau\, d\tau\, r dr$$
$$I_0\left(\frac{\gamma v_r p_T}{T}\right) K_0\left(\frac{\gamma M_T}{T}\right). \tag{18}$$

We would also like to point out that we consider the soft dilepton production in the low invariant mass region ($M < m_\rho$). There are many sources of dilepton in this mass range. If quark-gluon plasma is formed, there will be $q\bar{q}$ annihilation and $qq(\bar{q})$ scattering with virtual bremsstrahlung. It has been shown in Ref. [22] that in the case of low mass pairs the perturbative α_s-correction can be larger than the $q\bar{q}$ annihilation spectra, which has been calculated in the Born approximation. So, we consider this $q\bar{q}$ annihilation spectra as a lower limit and compare our result with it.

The yield of soft photons produced from bremsstrahlung process, $ab \to cd\gamma$ in an equilibrating plasma can also be obtained from (12) as

$$\frac{dN}{d^2q_T dy} = \int \tau\, d\tau\, r\, dr\, d\phi\, d\eta\, \frac{T^6 g_{ab}}{16\pi^4}$$
$$\times \int_{z_{min}}^{\infty} dz \frac{\xi(z^2T^2, m_a^2, m_b^2)}{T^4}\, (1 \pm \lambda_c \tilde{n}_c)$$
$$\times \Big[\lambda_a\lambda_b\lambda_d \sum_{n=0} \left(\pm e^{\beta E_c}\right)^n \frac{\mathcal{K}_1\left(z(n+1)\right)}{n+1}$$
$$+ \lambda_a\lambda_b\left(1-\lambda_d\right)\mathcal{K}_1(z)\Big]\Phi(s, s_2, m_a^2, m_b^2)$$
$$\times \left(q_0\frac{d\sigma^{\gamma}_{ab\to cd}}{d^3q}\right), \tag{19}$$

where $z_{min} = (m_a + m_b)/T$. The cross-section for the emission of a soft real photon is

$$q_0\frac{d\sigma^{\gamma}_{ab\to cd}}{d^3q} = \frac{\alpha}{4\pi^2}\, \frac{\widehat{\sigma}(s)}{q_0^2}\ , \tag{20}$$

where $\widehat{\sigma}(s)$ is defined as before (17) with J^μ replaced by the real photon current [14].

The photon production due to annihilation ($q\bar{q} \to g\gamma$) and Compton ($q(\bar{q})g \to q(\bar{q})\gamma$) processes in QGP have already been studied in fair detail by number of authors [23–26], in which thermal masses have been used to shield the singularities in the scattering cross-section. It has been shown [23] that this is equivalent to using the resummation method of Braaten and Pisarski [27] to regulate the divergences of the QCD rates for these processes. In the following we shall use the result of Ref. [24] for a chemically equilibrating plasma for comparison:

$$\frac{dN}{d^2q_T dy} = \int \tau\, d\tau\, r\, dr\, d\phi\, d\eta\, \frac{5\alpha\alpha_s}{18\pi^4} T^2 e^{-E/T}$$
$$\times \Big[\lambda_q^2\lambda_g\pi^2\Big[\ln\left(\frac{4ET}{k_c^2}\right) - 1.42\Big]$$
$$+ 2\lambda_q\lambda_g\left(1-\lambda_q\right)\Big[1 - 2\gamma + 2\ln\left(\frac{4ET}{k_c^2}\right)\Big]$$
$$+ 2\lambda_q\lambda_q\left(1-\lambda_g\right)\Big[-2 - 2\gamma + 2\ln\left(\frac{4ET}{k_c^2}\right)\Big]\Big] \tag{21}$$

In order to make a comparison of our calculation with that of using HTL [17] at the two loops level, we would like to discuss the following points. As discussed in Ref. [17,28], the HTL calculation leads also to a separation in two factors: the amplitude square of the scattering process without photon emission, and a factor which is nothing but the square of an electromagnetic current responsible for photon emission. Now one can easily convince oneself that this essentially amounts to the soft photon approximation we have used in our calculation. This can also be seen [17] quantatively when one compares the bremsstrahlung production of a soft virtual photon in the HTL calculation with that of the semi-classical approximation by Cleymans et al [29]. The semi-classical result agrees with the HTL result in its functional dependence but over-estimates the rate of production by 50%. This difference appears to be due to several simplifying assumptions made in Ref. [29]: (1) the photon momentum compared to the momenta of the constituents in the plasma has been neglected, but no phase space correction factor has been included. It has been discussed in Ref. [12,14] that the result grossly over-estimates without

the phase space correction; (2) the approximate electromagnetic current has been used; (3) also simply a factor unity in the exit channel has been assigned instead of considering the appropriate statistical factors. With proper inclusion, like in the present calculation, of these factors should lead to an agreement with the thermal field theory result. However, the contribution obtained in HTL using two loops is rather large compared to that in one loop [18,30], though both are of the same order in the coupling constant.

The real photon production in a QGP within the two loops using HTL [17] are also of the same order in the coupling constant as that of one-loop but there is quantative difference corresponding to different physical processes. The bremsstrahlung process in two loops gives a contribution which is similar in magnitude to the Compton and annihilation contributions evaluated up to the order of one loop [23]. In two loops there is also an entirely new mechanism for the production of hard photon through the annihilation of an off-mass shell quark and antiquark, where the off-mass shell quark is a product of the scattering with another quark or gluon. The contribution from this process completely dominates the emission of hard photons as energy increases, and it is higher by more than one order of magnitude for large energies [31,32]. In Ref. [32] it has also been shown that at the energy regime in which we are interested the bremsstrahlung from soft approximation is little higher than that from two loops, but the contribution from soft approximation falls off rapidly with increasing energy.

IV. RESULTS AND DISCUSSION

The parton fugacities for an expanding plasma are displayed in Figs.(1-2) in Sec.II with initial conditions obtained from SSPC model [3]. One can see that our results for the nonequilibrium soft photon (19) and low mass dileptons (15) due to bremsstrahlung differ from that of equilibrium rates. We would like to point out here that we cannot make a comparison of nonequilibrium results with that of equilibrium one as the present study is intended for RHIC and LHC energies. The SSPC model [3] predicts that the QGP likely to be formed in this energies will be far away from chemical equilibrium. In view of this, comparing the nonequilibrium results with that of equilibrium one will also be far away from reality. However, a close inspection suggests that our results are, indeed, quite different.

In Fig. 3, we give our results for the transverse mass distribution of low mass dileptons due to the bremsstrahlung process (solid curves) for $M = 0.1$ GeV, 0.3 GeV and 0.5 GeV, respectively, at both RHIC and LHC energies for an equilibrating plasma. We have also given the contribution from quark-antiquark annihilation process (dashed curve) for a comparison. We see that the annihilation contribution is essentially the same for $M = 0.1$ GeV, 0.3 GeV, and 0.5 GeV, thus representing a single curve in this figure. This is due to the fact that at M below the ρ mass the transverse mass distribution (see Eq.(18)) has a very weak dependence on M. But on the other hand it has a substantive dependence on the transverse mass M_T and thus shows M_T scaling [33,34]. This scaling is broken when we take into account the bremsstrahlung contribution from the QGP sector. We now see that in the region $M \leq 0.3$ GeV, the bremsstrahlung contribution out-shine the annihilation contribution at all M_T, while at $M \geq 0.3$ GeV, the distribution is dominated by the annihilation contribution with increasing available energy and invariant mass. This is because of the fact that quark fugacities are seen to lag behind the gluon fugacities at all time and all radial distances. Since our entire study is based on the description where the plasma as initially produced is gluon rich, the yield would be dominated by the quark-gluon bremsstrahlung. In this context, we see that the relative importance of bremsstrahlung contribution increases quite considerably for a chemically non-equilibrium plasma.

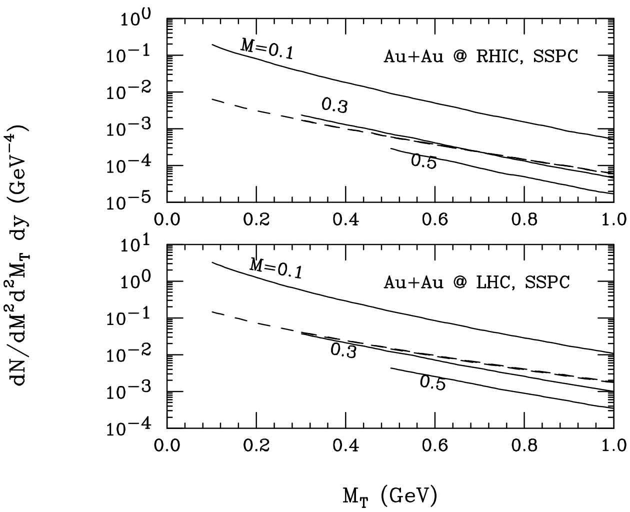

FIG. 3. The transverse mass distribution of low mass dileptons calculated with quark driven bremsstrahlung (solid curve) and quark- anti-quark annihilation (dashed curve; $M_T \geq M$) for $M = 0.1$ GeV, 0.3 GeV, and 0.5 GeV respectively at both RHIC and LHC energies.

We also plot the invariant mass spectrum of dileptons produced from bremsstrahlung (solid curve) and annihilation (dashed curve) processes in fig. 4. At $M \leq 0.3$ GeV, the bremsstrahlung process is seen to play the dominant role at both RHIC and LHC energies. We find that the low mass dilepton yield increases by a factor of 20 as we go from RHIC to LHC energies.

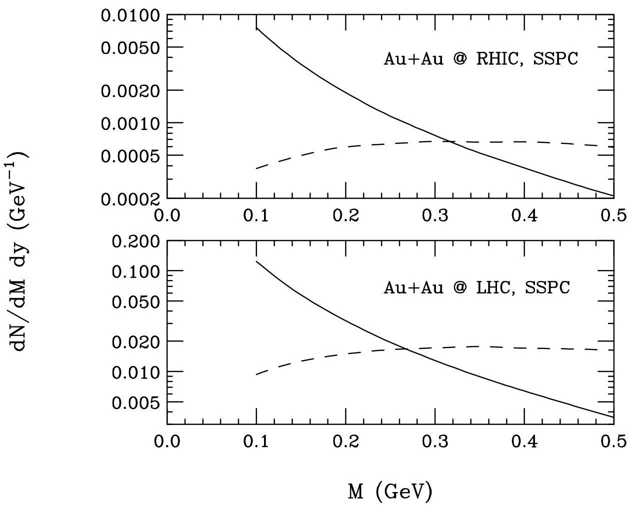

FIG. 4. The invariant mass distribution of low mass dileptons at both RHIC and LHC energies. The solid line represents the contribution from bremsstrahlung process in quark phase and dashed line represents the annihilation contribution from quark phase.

Figure 5 shows the transverse momentum distribution of soft photons at RHIC and LHC energies for an equilibrating plasma. We find that the contribution of the bremsstrahlung process (solid curve) dominates over that of Compton and annihilation process (dashed curve) up to $p_T \leq 0.5$ GeV at RHIC energy and at LHC energy, this dominance exists for $p_T \leq 0.6$ GeV, after which they fall rapidly. We envisage an increase by a factor of almost 10 in the soft-photon yield as we move from RHIC to LHC energies. For the low p_T regime, the photon rate has a substantial contribution from bremsstrahlung process whereas in the high p_T regime it is due to the hard processes (Compton and annihilation). As discussed already it will really be very interesting to study the production of single photons from QGP using the two loops results of Aurenche et al [17] along with the one-loop contribution in a chemically equilibrating plasma, since it has been suggested [3] that a chemically equilibrated plasma is likely to be produced in RHIC and LHC energies with a fairly large temperature. This indicates that QGP likely to be produced at those energies will have a larger life time, and most likely the photons from the quark matter with the new rates [17] may out-shine the photons from hadronic matter.

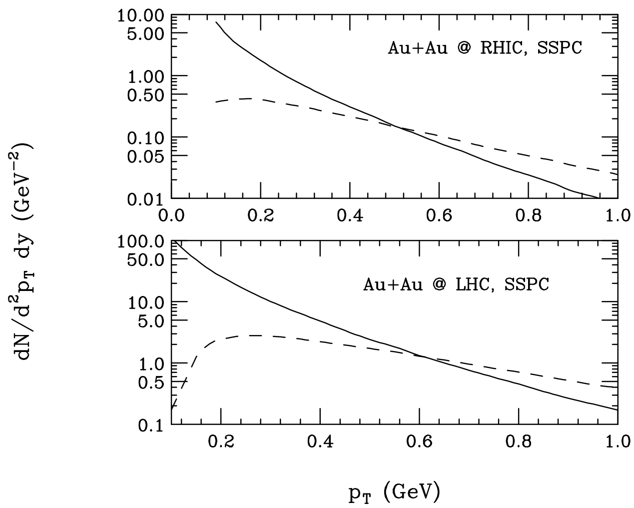

FIG. 5. The transverse momentum distribution of soft photons at both RHIC and LHC energies.

V. SUMMARY

In summary, we have studied low mass dileptons and soft photons production due to bremsstrahlung process at both RHIC and LHC energies from an equilibrating and transversely expanding quark- gluon plasma. We have chosen initial conditions from SSPC model in which the system is assumed to be in kinetic equilibrium at the proper time $\tau_i = 0.25$ fm/c but far away from chemical equilibrium [36]. In the present study, we restrict ourselves to photons and dileptons having energies larger than 100 MeV so that the Landau Pomeranchuk suppression [37] may not be substantial there. For dileptons, the bremsstrahlung contribution is seen to dominate up to $M \leq 0.3$ GeV and for photons this dominance exists for $p_T \leq 0.5$ GeV. We also find an increase by a factor of 15-20 in the low mass dilepton and soft photon yield as we go from RHIC to LHC energies where the life-time of plasma phase is expected to be large.

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